THE HEATING OF THE SOLAR CORONA

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Abstract

In this paper a discussion is given of the present state of the theory of the heating of the solar corona by shock waves. Arguments are presented why the main contribution to the mechanical energy flux is of acoustic origin, while estimates for the amount of acoustic energy generated in the convection zone as well as the deviations from isotropy are given. During propagation through the atmosphere acoustic waves develop into shock waves after a distance of a few scale heights in the chromosphere. The heating of the outer layers by dissipation of shock waves is found to be sufficient to account for the observed radiative and corpuscular energy losses.

Much emphasis is laid on the competitive role played by the four fundamental processes of energy transfer: mechanical heating, radiation, heat conduction and convection of energy in establishing the equilibrium structure of the corona. The atmosphere may be devided in several regions according to the predominance of one of the energy processes mentioned above.

The physical properties of the chromosphere and the solar wind are discussed only where they are intimately connected with the problem of the heating of the corona.

The most important aspects of the influence of a magnetic field on the structure and the heating of the corona in magnetically active regions are briefly mentioned. Special attention is paid to the strong chanelling of heat flow along the field lines and its consequences for the structure and dynamics of the chromosphere-corona transition layer.

1. Introduction

The theory of the heating of the solar corona has recently come to a stage where there appears to be reasonable agreement as to the fundamental energy processes concerned. The details of the energy balance are still a matter of discussion however. This is not surprising as the structure and the physical state of the corona are insufficiently known for making a detailed comparison between observations and theory possible. Moreover the distribution of the temperature and the chemical composition throughout the chromosphere and the corona, which is of vital importance to decide about the fundamental mode of energy transfer, is badly known. The structure and the physical properties of the chromosphere and the corona will not be discussed here as they have been treated elsewhere (van de Hulst, 1953; Thomas and Athay, 1961; Pagel, 1964; Shklovskii, 1965; Athay, 1965; Billings, 1966; Newkirk, 1967).

However, we must say a few words about the corona temperature. The corona is a gas with a very high temperature which became first apparent when Lyot (1932, 1937), with the coronograph, measured the width of the green coronal line and suggested that the broadening may be caused by thermal motions. Later Edlén (1943) found strong indications for a very high temperature after he identified the coronal lines as forbidden emission lines of highly ionized iron. Since then several attempts have been made to explain the high coronal temperature.

Early determinations of the temperature of the corona based on the theory of the ionization equilibrium yielded values in between $T = 5.10^5\ ^oK$ and $T = 10^6\ ^oK$. From the broadening of the coronal lines a value almost twice as large is found. This discrepancy was solved when Burgess (1964) showed that the process of dielectronic recombination, which has the effect of shifting the ionization equilibrium towards higher temperatures, has been overlooked. We may now

safely adopt a maximum coronal temperature in between $T = 1.5 \times 10^6$ ^oK and $T = 2.0 \times 10^6$ ^oK, with local fluctuations depending on the structure of the corona.

The hot corona overlying the low temperature solar photosphere is bound on the other side by cool interstellar matter. It is virtually impossible to deposit energy in the layers near the temperature maximum by means of radiative transfer and heat conduction only, since these processes essentially depend on the presence of temperature gradients. Neither does electrical energy seem to be a candidate for the heating of the corona. The coefficient of electrical conductivity is so large that Joule heating is insufficient for the energy supply.

The only process that can be taken seriously is the dissipation of mechanical energy. Dissipation of mechanical energy means that energy in ordered macroscopic motions is transferred into microscopic motions through viscosity or other collective processes. The ordered motions may be due to either a stream of particles (Thomas, 1948; Parker, 1958) or mechanical waves.

Another basic question to be settled is whether the corona is heated from the inside or from the outside.

Early theories advocated that the corona is heated by the accretion of interstellar matter (Bondi, Hoyle and Littleton, 1947; Biermann and ten Bruggencate, 1947; Pikel'ner, 1950). However, the amount of heat supplied by the accreted particles does not seem to be sufficient and the level of maximum temperature would be located at much too great a distance above the limb. A strong point in favour of the heating from the inside is the close relation of the spatial structure of the corona to the underlying chromospheric features. Accretion of interstellar particles may play some role in the energy-budget of the outer corona but it is certainly not important for the inner corona.

Actually the accretion theory in its present form should be revised regarding modern concepts of the physics of the outer solar

layers. It is therefore indicated to concentrate here our attention to the theory of the heating by mechanical waves having their origin inside the corona, notwithstanding the very few observational indications on which this theory is based; such as granulation, photospheric velocity fields, chromospheric oscillations and the increase of "turbulence" with the distance above the photosphere (Evans and Michard, 1962; Leighton, 1964). It was first suggested by Biermann (1946, 1948) and Schwarzschild (1948) that the outer solar atmosphere receives the heat input from compressional waves generated by the granulation, which was supposed to act as a field of pistons. It was also understood that the wave amplitude will increase as the waves propagate outward in the direction of decreasing density and that the waves will therefore steepen into a sequence of shocks. It thus became clear that within the framework of this theory the outer solar atmosphere could be heated by dissipation processes in shock waves (Schatzman, 1949; Schirmer, 1950).

In a discussion of the energy balance in the corona the first quantity to be evaluated is the amount of heat required to maintain the corona. Of particular importance is the ratio between the chromospheric heat input and what is left over to heat the corona. Two processes transport energy from the corona into interstellar space and may therefore be considered as energy losses, viz. electromagnetic radiation and corpuscular radiation or solar wind. However, in the inner corona heat conduction dominates the energy balance but the solar wind gradually takes over in the outer corona. Within a broad range of temperatures radiative losses are approximately independent of temperature and depend only on the particle density (see section 4, equation 12). If r_o is the radial distance of the "base" of the corona measured from the center of the sun, the total flux of electromagnetic radiation which has to be replenished from sources underneath is approximately given by

$$F_r = \int_{r_o}^{\infty} q(r) \left(\frac{r}{r_o}\right)^2 dr \approx \frac{RTq(r_o)}{2\mu g(r_o)} = 5.4 \times 10^4 \text{ ergs/cm}^2\text{sec} \quad (1)$$

where $q(r)$ is the radiative loss function, T is the temperature ($T = 1.5 \times 10^6$ °K), μ is the mean molecular weight, R is the gas constant and $g = GM_\odot/r^2$ is the acceleration of gravity. For reasons of simplicity it is assumed that the corona is isothermal. Most of the emission comes from the low coronal layers where the geometry is almost plane. The electron density at the base of the corona has been taken equal to $n_e = 10^9$ cm^{-3}. The flux of energy necessary to maintain the solar wind is approximately given by

$$F_s = \rho_E v_E \left(\tfrac{1}{2} v_E^2 + GM_\odot/r_\odot \right) \left(r_E/r_\odot \right)^2 \quad (2)$$

where ρ_E is the mass density and v_E is the velocity of the solar wind both taken at the earth orbit, r_E is the distance between the sun and the earth, $r_\odot$ is the radius of the sun and $GM_\odot/r_\odot$ is the gravitational potential energy of one gram of solar material. With $v_E = 5.10^7$ cm/sec and $n_E = 4$ cm^{-3} we obtain $F_s = 4 \times 10^4$ ergs/cm^{2}sec. Both the radiative and the corpuscular energy losses of the corona are of the same order.

From the level of maximum coronal temperature energy is conducted back to the chromosphere through the transition region and outward into the interplanetary medium. The latter flux of energy is somewhat smaller than the solar wind energy flux. It seems that the greater part of this flux is used to drive the solar wind beyond the levels where the dissipation of mechanical energy is the main heat source. This would mean that the outward conductive energy flux can not be considered as a real energy loss, independent of the convective energy loss (Noble and Scarf, 1963, 1965; Kopp, 1968).

Altogether it can be concluded that a flux $F_m \approx 10^5$ ergs/cm^{2}sec in the form of waves should reach the base of the corona.

The chromosphere only looses energy by radiation essentially in the form of H and H^- continuum, Balmer lines and resonance lines of

elements like Ca, Si, Mg.

Osterbrock (1961) has estimated the chromospheric radiative losses. He gives a value for the radiative flux $F_{chrom} = 4 \times 10^7$ erg/cm^2 sec, most of it being H^- emission. This estimate has been analysed by Dubov (1965), who argued that the H^- continuum is emitted in a shallow layer of a few hundred kilometer thickness located in the lowermost part of the chromosphere. He obtained a value of the order $F(H^-) \approx 10^6$ erg/cm^2 sec for the H^- emission and concluded that less than 5.10^6 erg/cm^2 sec is emitted predominantly in the Balmer lines.

Recently, Athay (1966a, b) reviewed the most important sources for chromospheric and coronal energy losses in order to get an observational parameter to compare theories of chromospheric heating.

He again found that the H^- flux is about an order of magnitude smaller that has been estimated previously by Osterbrock. He also found that the emission originates in a fairly thin layer. Athay's results are summarized in table 1.

Table 1

Chromospheric Radiative losses (after Athay)

height	0 - 500 km	500 - 1000 km	1500 km	> 1500 km
radiator	H^-	Balmer series	Lyman α	UV metal lines
scale height	100 km	500 km	1000 km	10^4 km
radiative flux (erg/cm^2 sec)	4×10^6	10^6	3×10^5	3×10^5

The total radiative flux leaving the chromosphere is therefore $F_{chrom} = 5.6 \times 10^6$ erg/cm^2 sec.

From these non-thermal radiative energy losses given above Biermann (1966) has estimated the shock strength and the energy density of the shock waves which are supposed to maintain the chromosphere.

Much more accurate estimates of the chromospheric energy losses are urgently needed in order

to determine which heating mechanism is working at different heights.

Much of the hypothetical nature of the theory of the corona would vanish if better chromospheric data were available.

A discussion of the heating by waves naturally falls into three classes of problems; the generation, the propagation and the dissipation of the waves.

The generation of mechanical waves in the sun will be discussed in the subsequent section.

2. The Generation of Waves in a turbulent atmosphere

The layers directly underneath the chromosphere can be devided into stable layers and unstable layers.
The stable layers are essentially in radiative equilibrium. The unstable layers are in convective equilibrium, which means that most of the energy is transported from the inside by convective motions. There is no sharp separation between the stable and the unstable layers.
In the upper convection zone radiation is gradually taking over the energy transport, while on the other hand convective motions penetrate into the stable layers (Moore, 1967). The upper convective layers and the transition layers between the unstable and the stable solar atmosphere which are both in highly turbulent motion are supposed to be the breeding grounds for the mechanical waves.

Under the influence of gravitational forces and a magnetic field three fundamental modes of wavepropagation appear in a compressible medium which can be distinguished by the way the divergence and the vorticity (curl $\vec{v}$) of the velocity field are propagated (Lighthill, 1967). Ordinary sound waves depend on compressibility alone, and they can propagate in all directions.
The wavenumber surface or slowness surface is a sphere with a radius given by $\vec{k}^2 = \omega^2/c^2$, where $\vec{k}$ is the wavevector, ω is the frequency and c is the sound velocity. The group velocity is in the same direction as the phase velocity. In a gravitational field $\vec{g} = (0,0,-g)$, (the z-direction being directed radially outward) and ignoring compressibility, internal gravity waves can propagate for all frequencies $\omega < \omega_2$, where ω_2 is the Väisälä-Brunt frequency given by

$$\omega_2^2 = -\frac{g}{\rho_o}\frac{d\rho_o}{dz} - \frac{g}{c^2} \qquad (3)$$

The wavenumber surface is a cone with semiangle given by $\sin\theta = \omega/\omega_2$. In this case the group velocity is perpendicular to the phase velocity. The divergence of the velocity is zero and the vertical component of the vorticity is constant.
It should be noted here that the atmosphere is convectively unstable if $\omega_2^2 < 0$. As a consequence internal gravity waves cannot be generated in the convective layers of the sun. The third type of wave is the Alfvén wave, which propagates along the magnetic lines of force in a magnetoplasma. Again the divergence of the velocity is zero. The vorticity propagates along the field lines. For Alfvén waves the wavenumber surface consists of two planes perpendicular to the direction of

the magnetic field given by $(\vec{B}.\vec{k}) = \pm\ \omega\ (4\pi\rho)^{\frac{1}{2}}$.

Sound waves are modified by gravity. The wavenumber surface is transformed into an oblate ellipsoid and the group velocity being directed perpendicular to the wavenumber surface obtains therefore a preference for the vertical direction (figure 1). The oblateness increases as the frequency decreases until the critical frequency $\omega_1 = (\gamma g)/(2c)$ is reached, below which acoustic wave propagation is impossible. In their turn gravity waves are modified by the compressibility so that the group velocity is no longer perpendicular to the phase velocity. The wave number surface becomes a hyperboloid with the original wave cone in an incompressible medium as the asymptotic cone approached at large wave numbers. Neither sound waves nor gravity waves can propagate in the frequency range $\omega_2 < \omega < \omega_1$ (Whitney, 1957; Eckart, 1960; Hines, 1960; Tolstoy, 1963; Whitaker, 1963; Moore and Spiegel, 1964; Souffrin, 1966; Stein, 1967).

In a compressible plasma in a magnetic field the well known fast-mode, slow-mode and Alfvén-mode m.h.d. waves can propagate (van de Hulst, 1949; Osterbrock, 1961; Lighthill, 1967).

Under the combined action of non-parallel gravity and magnetic fields an interesting mode coupling between gravity waves and Alfvén waves is possible, and no longer any frequency limitations for propagation exists. The propagation surface approaches that for pure Alfvén waves when $\omega > \omega_2$ and the group velocity approaches the direction of the magnetic field (Lighthill, 1967). It is important that it appears possible in this way for a gravity wave to transform into an Alfvén wave.

The next question is where these wave modes can be generated and how much energy is fed into each of them. The generation of waves is of course not independent of their propagation. For example it is unlikely that at some place much energy can be fed into wave modes whose propagation velocity is small compared to the propagation velocity of locally important other disturbances in the medium.

For instance we do not expect that Alfvén waves are generated in the turbulent layers of the solar atmosphere because the Alfvén velocity is small compared with the sound velocity. Only in active regions they may contribute to the non radiative energy flux in the photosphere.

For example if the magnetic field strength is 2 gauss the Joule-heating damping length is of the order of 100 meter in the upper convection zone owing essentially to the smallness of the Alfvén velocity as compared to the sound velocity (Osterbrock, 1961). The layers which are located a few hundred kilometers higher are conducive for the generation of Alfvén waves because the Alfvén velocity is comparable with the sound velocity.

On the other hand sound waves can be generated efficiently in the turbulent layers of the convection zone as well as in the layers of turbulent penetration, while gravity waves can only be generated in the stable layers of turbulent penetration.

For the generation of sound waves the situation is as follows; in a homogeneous isotropic compressible turbulent atmosphere sound waves are generated by fluctuations in the momentumflux (Lighthill, 1952, 1954, 1960, 1962). For small Mach numbers there is no feedback from the generated sound upon the turbulence. Then the major contribution to the sound emission is quadrupole radiation.

The acoustic power generated by isotropic turbulence has been estimated by Proudman (1952) and is given by

$$P = 38 \frac{\rho_o \langle u^2 \rangle^{\frac{3}{2}}}{\ell} M^5, \qquad (4)$$

where ℓ is a characteristic length for the turbulence correlation, M is the turbulent Mach number. Stein (1967) gives a value which is an order of magnitude larger than Proudman's estimate, which may be understandable since the power output depends strongly on the shape of the turbulence spectrum as well as on the adopted turbulence correlations.

The flux of mechanical energy, which is the acoustic power integrated over depth has been estimated by De Jager and Kuperus (1961), Osterbrock (1961) and Kuperus (1965) on the basis of expression (5) using available models of the hydrogen convection zone.

Figure 2 shows the acoustic power as a function of the geometrical depth measured from the level where the optical depth in the continuum $\bar{\tau} = 1$ on the basis of Proudman's expression (5) using a model for the convection zone given by Vitense (1953).

It is striking that all sound is essentially generated in a thin layer of about 60 km thickness located at the top of the convection zone. Owing to the strong dependence of the acoustic power on the velocity

for which one essentially depends on theoretical models, the values given for the flux of mechanical energy are at best order of magnitude estimates. Most estimates result in a value around $F_m \approx 3 \times 10^7$ ergs/cm^{2}sec.

There are many effects that may influence the amount of energy generated in the sound waves.

First the waves are probably emitted in bursts and are not monochromatic. This means that the turbulent velocity has to be compared with the group velocity of the sound waves in a gravity field. Especially close to the critical frequency ω, this results in a large increase of the "effective" Mach number of the turbulent motions. It is expected that due to this effect most sound is radiated close to the critical frequency, but this has te be properly worked out. Actually Kato (1966) found that the field of acoustic waves generated by turbulence in an atmosphere can be strongly peaked close to the critical frequency ω_1.

Secondly some increase in the power output may be gained because of additional dipole radiation with a preferred direction due to forcing terms like a magnetic field or the gravity field. This effect has been considered for a magnetic field by Osterbrock (1961) and for a gravity field by Unno and Kato (1962) and Unno (1964).

A third effect is due to the "relativistic" enlargement of the sound generating turbulent eddies, which are embedded in the mean convective flow. The quadrupole radiation intensity in a direction making an angle θ to the direction of the mean flow velocity v_c is given by the equation 5 multiplied by a factor $(1-M_c\cos\theta)^{-5}$, where $M_c = v_c/c$ Ffowcs Williams (1963). Figure 3 shows the deviations from isotropy for different values of M_c. This means that on the sun more sound is generated in the upwelling region than in the downward moving regions (Kuperus, 1965).

However, if sound waves propagating downward are refracted upward by the temperature gradient, approximately the same intensity of the sound is expected in the upwelling as well as in the downward moving regions. In the intermediate regions the sound intensity is then lower. Finally, the presence of turbulent magnetic fields results in an enhancement of the sound production (Kulsrud, 1955). The effect is only of importance when the turbulent magnetic energy is comparable with the turbulent kinetic energy. In that case we do expect the magnetic fields to extract so much energy out of the turbulent motions that it does not seem fair to compare the power output with and without a turbulent magnetic field using the same velocity distribution.

Since there is no satisfactory theory of m.h.d. turbulence this problem cannot be solved at this moment. Sofar the generation of sound waves.

As far as gravity waves are concerned it has only been suggested by Whitaker (1963) that the solar corona could be maintained because of the dissipation of gravity waves by thermal conduction. However, it is difficult to estimate the flux of gravity waves because we do not know the extension of the layers of turbulent penetration, where gravity waves can be excited. Stein (1967) derived an expression for the flux of gravity waves being proportional with $(\ell_o/h)^5$, where ℓ_o is a characteristic scale of the turbulence and H is the local scale height of the atmosphere. From his analysis it appears that the energy radiated by gravity waves is so large that it would undoubtedly result in a strong damping of the turbulent motions so that the extension of the penetrative turbulence in the stable layers is reduced.

On the other hand radiative dissipation in the optically thin stable photospheric layers strongly reduces the efficiency of the generation of gravity waves because the radiative relaxation time in the low chromosphere is smaller than the Brunt - Väisälä frequency. This causes a strong damping of gravity waves (Souffrin, 1966; Schatzmann and Souffrin, 1967).

An other argument against gravity waves to carry appreciable amounts of mechanical energy to the outer layers is their strong reflection in the photosphere - chromosphere temperature minimum. The transmissivity of gravity waves is a few orders of magnitude smaller than the transmissivity of acoustic waves (Uchida, 1965, 1967). Outside the active regions and the chromospheric network Alfvén waves are not likely to contribute to the energybalance of the outer atmospheric layers. Gravity waves may be generated and there is a possibility that part of the gravity wave energy is fed into Alfvén waves especially in the magnetically active regions.

Concluding it appears as if the dominant contribution to the flux of mechanical energy is due to ordinary sound waves.

3. The formation and the dissipation of shock waves.

Acoustic waves are weakly damped as long as the velocity amplitude is small compared with the local sound velocity. They do not contribute substantially to the heating of the photosphere and the chromosphere unless their amplitude has become comparable with the sound velocity. However, even if there were no attenuation by dissipative processes the flux of mechanical energy is decreased because of the refraction and reflection of acoustic waves in a medium with the temperature increasing outward (Schatzman, 1949).

The reduction is greater if the energy is transported by fast mode magnetohydrodynamic waves in a magnetic plasma (Osterbrock, 1961). Kopp (1968) compared the transmission coefficient for acoustic waves incident on the interface of two media with different temperatures, taking the partial reflection of waves inside the so called critical cone into account. For the interface between a medium with a temperature of 7000 ^{o}K and a medium with a temperature of 1.4×10^{6} ^{o}K, which schematically represents the transition between the chromosphere and the corona, he finds for the fraction of the total acoustic radiation which is incident isotropically from below a value of 4.4×10^{-4}. Actually this value is an upper limit since dissipation has been ignored. However, there are two basic arguments why the transmissivity of the solar atmosphere for acoustic waves is much larger. First the emission of sound waves is preferentially in the vertical direction, particularly for waves with a frequency close to the critical frequency and as had been pointed out in section 2 sound waves generated in convective turbulence are predominantly emitted in the direction of the convective motion. Consequently the emission of sound in the upwelling and downward streaming regions of the solar convection zone is highly anisotropic (Kuperus, 1965). Secondly the transmission of sound waves is certainly larger than has been predicted by simple geometrical acoustics because at the time the

acoustic waves reach the chromosphere-corona interface they have been steepened into shock waves. Since the density decreases exponentially with height the velocity amplitude increases with height approximately as $v \approx \exp(z/2H)$, where H is the density scale height of the atmosphere. The average height of shock formation is a few hundred kilometer above the photosphere (Bird, 1964c; Schatzman, 1967). It is therefore not likely that dissipation of linear waves by thermal conduction might account for the heating of the chromosphere as has been suggested by Sturrock (1964) on the basis of a very uncertain temperature-density relation in the chromosphere. Numerical solutions of the full gasdynamic equations in a gravity field without radiation show that after a few scaleheights a shockwave with a sawtoothshape profile has been formed out of a sinusoidal wave (Kuperus, 1969). An example of such a calculation for a wave with a period of 200 sec. is shown in figure 4. It has been shown that the critical angle under which waves may penetrate the interface depends on the Mach number of the shock waves (Kopp, 1968). Hence, the transmission coefficient for strong shocks incident obliquely on a temperature discontinuity is larger than for weak shocks or acoustic waves.

During the propagation of shock waves through the atmosphere some of the shock wave energy is transferred into thermal motions due to viscous processes. The dissipation of shock waves is an immediate consequence of the increase in entropy of the gas behind the shock-front due to irreversible processes like viscosity and heat conduction in the transition layer. The entropy distribution through the shock-front can be expressed in terms of the viscous heating and the cooling or heating of the fluid by heat conduction and radiative transfer. If μ is the kinematic viscosity, $\varkappa$ the coefficient of thermal conduction and F the radiative flux, the increase of entropy $\frac{DS}{Dt}$ is given by

$$\rho T \frac{DS}{Dt} = \frac{4}{3} \mu \left(\frac{\partial u}{\partial x}\right)^2 + \frac{\partial}{\partial x} \left(\varkappa \frac{\partial T}{\partial x} \right) + \frac{\partial F}{\partial x} , \qquad (5)$$

where T is the temperature and u the material velocity (Zeldovich and Raizer, 1966).

The exact way to evaluate the heating of a fluid element when it is compressed by a shock wave is by means of equation 6. Sofar this method has only been used by Mäckle (1968) for the determination of an accurate heating function of the chromosphere.

However, if one is only interested in the final state of the gas after the shock has passed and not so much in the distribution throughout the front the method described above is unnecessarily complicated, because the jump in entropy is uniquely determined by the shock strength. Actually no knowledge of the physical processes, like viscosity and heat conduction, in the transition layer is required to determine the final amount of heating.

An exception should be made for an optically thin medium where the photon mean free path is much larger than the radiative shock width. Otherwise dissipative processes only determine the width of the shock and not the jump conditions.

In a quasi-steady atmosphere permeated by a sequence of shock waves separated by the accompanying system of expansion waves the entropy rise after each shock is balanced by non adiabatic processes like heat conduction and radiation during the expansion.

The periodic process of shock compression can be divided into three different processes, as is illustrated schematically in figure 5. After a sudden compression (0-1) the entropy increases by an amount ΔS. In a p-V diagram the state of the gas moves along the **shock**-adiabatic from the adiabatic S_0 to the adiabatic S_1. The subsequent expansion essentially exists of a slow isentropic expansion, which reduces the pressure to its original value p_0 and an isobaric cooling by radiation losses and heat conduction (Schatzman, 1949). Actually both processes occur simultaneously. For weak shocks the expansion is practically over before the cooling becomes appreciable. For strong shock waves radiative cooling just behind the shock (state 1) may be so important that it is more appropriate to consider the cooling process before the adiabatic expansion (Weymann, 1960).

The real dissipation during this quasistatic cyclic process is $\oint p dV$ which is indicated in figure 5 by the hatched surface for the socalled Schatzman path and by the dotted surface for the Weymann path. The amount of dissipation is immediately connected to the net increase in enthalpy after the medium has been expanded to its initial preshock conditions. For weak shock waves the enthalpy rise is given

by

$$\Delta H = \frac{2}{3} a_o^2 \frac{(M^2-1)^3}{(\gamma+1)^2} , \qquad (6)$$

which is of the third power in (M^2-1), where $M = V/a_o$ is the Mach number of the shock waves and a_o is the sound velocity in the medium in front of the shock (Schatzman, 1967). Due to the dissipation shock waves are damped during their propagation. To find the damping of shock waves the similarity hypothesis of Brinkley and Kirkwood, (1947) stating that the expansion profile remains invariant during the propagation has been frequently applied (Schatzman, 1949; Weymann, 1960; Ono, Sakashita and Ohyama, 1961; Osterbrock, 1961; de Jager and Kuperus, 1961; Dubov, 1960a, b; Saito, 1964; Ulmschneider, 1967). As a matter of fact the damping and the propagation properties of shock waves in the solar atmosphere are interrelated since the shock wave Mach number is strongly influenced by the distribution of temperature and density. On the other hand these gradients are determined by the local energy balance in which the dissipation of shock waves is the fundamental process.

For reasons of simplicity some authors preferred to use an approximate transmission law for shock waves which was derived by splitting up the atmosphere in a few discrete layers (Chisnell, 1955; Osterbrock, 1961; Ono, Sakashita and Ohyama, 1961; Kuperus, 1965), while others tried to solve the propagation equations taking into account the effect of gradients (Bird, 1964a, b, 1965a, b; Ulmschneider, 1967; Kopp, 1968).

The theories of the last category make use of the method of Whitham, which is based on the theory of characteristics. Whitham (1958) calculated the propagation of shock waves in an inhomogeneous medium under the assumption that disturbances moving along the negative characteristics do not interact with disturbances moving along the positive characteristics. Bird (1964a) showed that Whitham's

method gives satisfactory results for strong shock waves. It should be mentioned here that in the outer coronal layers where heating is still necessary the ordinary viscosity is so small that other collective processes take over the role of the collisions and assure the entropy jump through the collisionless shocks. It has been suggested by d'Angelo (1968) that Landau damping of ion-acoustic waves is more important in the high coronal layers than normal viscous damping, while Barnes (1968) showed that the collisionless solar-wind plasma at the earth's orbit may be heated by fast mode magnetohydrodynamic waves.

4. The energy balance in the outer solar atmosphere.

The structure and the physical state of the outer solar atmosphere is entirely determined by the detailed energy balance at different heights. It has been argued that the basic energy source for chromospheric and coronal heating is the dissipation of mechanical energy by shock waves. In a static atmosphere the dissipated energy is transferred by radiation and heat conduction only, but when the atmosphere is in steady motion the energy equation in spherical symmetry is given by

$$\frac{1}{R^2}\frac{\partial}{\partial R}\left\{R^2\left(\rho u\left[\tfrac{1}{2}u^2+h-\frac{GM_\odot}{R}\right]\right)\right\} = q_h - q_r + \frac{1}{R^2}\frac{\partial}{\partial R}\left(R^2\varkappa\frac{\partial T}{\partial R}\right), \quad (7)$$

where the left hand side is the energy which is convected in the radial direction with the velocity u, h is the specific enthalpy, q_h is the heating per gram of matter, q_r is the loss of energy by radiative processes. The third term on the right hand side represents the energy loss by thermal conduction, $\varkappa$ is the coefficient of thermal conduction (Bird, 1964 a,b; Ulmschneider, 1967; Kopp, 1968). According to the way the energy is transported Kopp (1968) devides the outer solar atmosphere into seven distinct regions, which are shown in figure 6. Region I is the chromosphere where mechanical heating is balanced by radiation. Heat conduction is unimportant. The temperature structure is largely unknown mainly due to insufficient knowledge of the efficiency of the numerous radiative processes which take place in the chromosphere. The thickness is estimated of the order of 1000 km (Athay, 1966 b).

Region II is the lowermost part of the transition region between the chromosphere and the corona in which the temperature starts to rise suddenly. There thermal conduction plays a competitive role with radiation. In region III, which is also still part of the transition region radiation is less important than heat conduction. The conductive energy flux is almost constant, and hence largely

determines the temperature profile. In region IV the temperature increases from about 10^6 °K to 2 x 10^6 °K which is roughly the maximum temperature. Mechanical heating is very important at the temperature maximum, heat is conducted both into the outer corona as well as back into the chromosphere. Since the layers below the level of maximum temperature have in common that the energy which is convected outward by the solar wind is negligible in comparison with radiation and heat conduction these layers may be safely considered to be in hydrostatic equilibrium. This is certainly not true for the outermost layers. Hydrodynamic expansion becomes increasingly important in region V. The upper boundary of region V is located at a level where the shock waves are exhausted. Above this level in region VI neither mechanical heating nor radiative losses play a role in the energybudget. The temperature structure in region VI is entirely determined by the interplay of thermal conduction and hydrodynamic expansion, if no other unknown heat sources are present. At large distances from the sun heat conduction becomes negligible and the remaining energy is convected outward only by adiabatic hydrodynamic expansion of the solar corona (region VII).

The average distribution of the density and the temperature of a steady shock-heated atmosphere can be computed from the energy-equation (8) together with the momentum equation

$$u\frac{\partial u}{\partial R} = -\frac{1}{\rho}\frac{\partial p}{\partial R} - \frac{GM_\odot}{R^2} \qquad (8)$$

and the equation of continuity

$$-\frac{1}{R^2}\frac{\partial}{\partial R}(\rho u R^2) = 0 \qquad (9)$$

after having specified $\varkappa$, q_r, q_h and the boundary conditions. In a fully ionized plasma a reasonable approximation for $\varkappa$ is

$$\varkappa = 6 \times 10^{-7} \times T^{\frac{5}{2}} \qquad (10)$$

(Chapman, 1954).

The radiative losses of a unit volume of optical thin coronal matter consists of

a. free-free emission of hydrogen and helium.

b. free-bound continuum of hydrogen and helium.

c. line emission from hydrogen and helium.

d. permitted line emission from heavy elements.

e. forbidden line emission from highly ionized heavy elements.

Orrall and Zirker (1961) found that the total radiative losses are approximately independent of the temperature and are given by

$$q_r = 1.76 \times 10^{-23} n_e^2 \quad \text{ergs/cm}^3\text{sec.} \qquad (11)$$

Bodenheimer, Brandt and Robbins (1963) gave a similar expression

$$q_r = 2.30 \times 10^{-23} U\, n_e^2 \quad \text{ergs/cm}^3\text{sec.} \qquad (12)$$

where U is an uncertainty factor due to our poor knowledge in relative abundances and oscillator strengths.

The expression (12) was used by Kuperus (1965) and the expression (13) by Kopp (1968).

The heating per unit volume q_h due to the dissipation of shock waves is proportional to the increase in enthalpy during a shock and expansion cycle.

For weak shock waves ($M_s \approx 1$) we find from equation (7)

$$q_h = \frac{\rho}{P} \Delta h = \frac{16}{3} \frac{\rho c^2}{P(\gamma+1)^2} (M_s-1)^3, \qquad (13)$$

where P is the period of the shocks and M_s is the Mach number.

Kuperus (1965) used a slightly different expression for q_h based on the assumption that the wavelength of a sequence of shocks remains constant (Landau and Lifschitz, 1959).

In the limit of weak shocks q_h is again given by equation (14). The equations 8 - 14 can be solved as soon as M_s is known as a function of the radial distance R. The exact expression for $M_s(R)$ is a very complicated function which in general can be found from the implicit relation

$$\frac{1}{M_s}\frac{dM_s}{dR} = \Pi(M, M_s, T, R), \qquad (14)$$

where M is the Mach number of the average flow (i.e. the solar wind), given by the critical solution of the solar wind equation (Parker, 1963, 1965), and T is the temperature (Bird, 1964 a,b; 1965; Ulmschneider, 1967; Kopp, 1968). If one is interested in the global properties of the outer solar atmosphere the concept of an absorption coefficient $\varkappa_s$ is useful. If F is the flux of mechanical energy $\varkappa_s = F^{-1}\, dF/dR$ which is the inverse damping length. The relation between $\varkappa_s$ and q_h is given by $q_h = \varkappa_s F$, where $\varkappa_s$ depends implicitly upon F. A small variation of F or M_s or a slightly different boundary condition propagates through any numerical calculations and adds a great factor of uncertainty to any model which has been presented. It may occasionally result in a complete different type of solution. Yet a comparison of different models is instructive as it shows the effect of different assumptions and boundary conditions and leads to a better understanding of the physics of the solar corona.

5. A discussion of models of a steady heated outer solar atmosphere.

The early investigations of Schatzman (1949) and Schirmer (1950) did not result into a quantitative model of the outer solar atmosphere. By a simple balance of the shock wave dissipation (heating) against the radiative losses Schatzman concluded that the transition from chromosphere to corona should be very abrupt essentially because at low densities the gas is not able to radiate the energy deposited by the shock waves in a low temperature configuration. Thermal conduction was not taken into account. This explanation of the sudden rise of temperature from chromosphere to corona has originally been put forward by Shklovskii (1944). The same analysis was followed by Dubov (1960 a,b) who argued that the magnetic fields do not appreciably change the amount of heating and by Weymann (1960) who concluded that the transition region between chromosphere and corona may be characterized by one or more steep rises in temperature. It should be stressed here that the neglect of thermal conduction is not allowed, in contrast to that of hydrodynamic expansion, at least in those layers where hydrostatic equilibrium may be applied. The possible importance of thermal conduction in the solar corona was already advocated by Alfvén (1941). He pictured the solar corona as an extended atmosphere presumably heated by a stream of high energy particles emanating from the chromospheric layers. It was supposed that the particles thermalize close to the surface of the sun. From this level thermal conduction should transport the heat outwards as well as inwards. In many respects Alfvén's supposition came close to what seems generally accepted at the moment. Later Giovanelli (1949) pointed out that thermal conduction constitutes the main contribution to the energy transport in the lower layers. He thought that the mechanical energy was transported by convective motions in the photosphere and chromosphere and that the mode of energy transfer changes from convection to thermal conduction in

the transition region. Various models of the transition region and the outer corona based on thermal conduction only have been constructed (Woolley and Allen, 1950; Oster, 1956; Chapman, 1957; Unsöld, 1960).

Osterbrock (1961) computed the dissipation of shock waves in the chromosphere based on the Brinkley-Kirkwood theory including the effect of a magnetic field. He found that the dissipation was sufficiently large to heat the chromosphere but largely insufficient to cause the heating of the corona. His result is essentially derived from geometrical acoustics, which is only correct for small amplitude waves and for waves with wavelengths much smaller than the scale height. Both conditions are not fulfilled for shock waves in the solar atmosphere.

Besides, the emission of acoustic waves in the photosphere is not isotropic at all but occurs very much in the vertical direction, as has been argued in section 2. It seems that the arguments given by Osterbrock, that the corona should be heated by other wave modes can not be maintained. The first model of the outer solar atmosphere taking into account heat conduction as well as radiative losses was made by De Jager and Kuperus (1961) and later refined by Kuperus (1964, 1965). Four differential equations giving the derivatives of the density (from hydrostatic equilibrium), temperature (from the equation of heat conduction), mechanical energy flux (from the propagation of shock waves) and conductive energy flux (from the energy equation) were integrated downward starting in the corona at the level of maximum temperature where $dT/dz = 0$. The results of two integrations are given in figure 7 for comparison with other model calculations. From available chromospheric models and the particle density at the "top" of the chromosphere the temperature of the corona was found a little below $T = 10^6\ ^oK$.

A reduction of the thickness of the chromosphere naturally leads to

an increase in coronal temperature. The calculations were made using a shock wave period of 300 sec. The results strongly depend on the adopted period. No solar type corona was found for a period of 100 sec. The transition region is characterized by a steep temperature gradient around $T = 10^5$ °K. This is a natural consequence of the near constancy of the conductive energy flux in the lower levels of the transition region.

Uchida (1963) studied the effect of a magnetic field on the heating of the corona again taking heat conduction into account. It is interesting to note that he performed the integration starting from the low chromosphere. This was possible because he used a temperature dependent radiative loss function. Uchida found a higher coronal temperature than Kuperus did, probably due to the neglect of shock wave reflection. It is not clear how sensitive Uchida's solutions are to the boundary conditions at the base of the chromosphere.

A different approach to the equilibrium structure of the solar corona was made by Bird (1964 a,b), using the method of Whitham. He argued that for a large interval of shock wave periods a quasi-steady shock heated atmosphere adjusts in such a way that the shock strength is not seriously altered. Bird neglected heat losses by radiation and heat conduction. In his model the energy convected outward is balanced by the shock dissipation. The hypothesis of a constant shock strength replaces the energy equation. This is a great simplification of the problem. The temperature profile of the atmosphere is chosen such as to assure the constancy of the shock strength. This does not seem unreasonable for weak shocks as long as heat losses by conduction and radiation may be neglected.

Later Bird (1965) used an energy equation with the proper heat function but still neglecting heat conduction and radiative losses to check his former computations. The integrations were performed outwards starting with a Mach number of the average flow which matches the critical solar wind solution. It was found

that the shock Mach number was not constant during the propagation. Yet the overall behaviour of the temperature profile was almost similar to the one based on a constant shock strength. The maximum temperature of the corona changes only by a few percent around $T = 1.6 \times 10^6\ ^oK$, while increasing the initial heat input through several orders of magnitude. Apparently at a certain temperature almost all the extra heat is used for the hydrodynamic expansion. Two of Bird's model calculations for different shock wave Mach numbers M_s are given in figure 7. It should be stressed at this point that the omission of heat conduction and radiative losses will never yield a realistic solution for the transition layer and inner corona. The difference between Bird's calculations and the other calculations which include heat conduction is striking especially in the transition layer. A major feature of Bird's models is that the maximum coronal temperature is reached at a distance of about one solar radius above the limb. This does not seem to be at all in accordance with observations. However, Bird's analysis has been very stimulating as it was the first quantitative attempt to connect the shock wave heating theory of the corona with the theory of the solar wind.

The full hydrodynamic equations with radiative losses and heat conduction taken into account were solved by Ulmschneider (1967). He also included the effect of a variable $\gamma = c_p/c_v$. His solutions are very sensitive to the Mach numbers of the average flow at the base of the atmosphere (M_o). If M_o exceeds a certain value, a physically unrealistic solution with an ever outwardly increasing temperature is found. At smaller values of M_o it is possible however to find a critical solar wind solution where the total energy convected outwards is larger than the total radiative loss. For still smaller values of M_o the convective energy loss becomes smaller than the total radiative loss. Ulmschneider rejects these latter solutions on the basis of the very steep temperature gradients

involved. His critical solution is given in figure 7. The maximum temperature is higher than that obtained from previous model calculations. Heating by shock waves has to take place unto a height of at least h = 70000 km, this is the level where the maximum temperature is reached. The Mach numbers M_s and M are treated in this theory as independent parameters. However, the Mach number of the average flow at the base of the atmosphere, M_o, must be dependent on the shock wave Mach number, M_{so}, at the base. It seems that a more general treatment with one parameter less should be possible.

An extended critical investigation of the possibilities and limitations of the theory of the equilibrium structure of a shock heated corona was made by Kopp (1968). Using shock expansion theory he found that the damping of shock waves by their wakes is a first order effect. The neglect of this effect causes apparent excessive heating of the outer corona. With this stronger damping and heat conduction taken into account a much steeper temperature gradient is found, and as a consequence the maximum temperature is again reached much closer to the solar surface than in Bird's analysis. A model for the outer solar atmosphere was obtained by integrating inwards the hydrodynamic equations for shock propagation specifying boundary conditions at three points, namely at the base, in the critical point and at infinity. One of his models is given in figure 7. Kopp showed that the observed density distribution is a rather poor means to inform about the nature of coronal heating. He found that for a wide range of conditions varying from a hydrostatic, isothermal corona to an expanding, conducting corona the computations can be made as to match the observed density distribution as far as $r = 6\ r_\odot$.

6. The influence of magnetic fields.

Most theories of the structure of a shock heated corona do not take into account the numerous modifications which occur in the presence of a magnetic field. These render it dangerous to draw general conclusions based on qualitative arguments only. However, it is possible to consider special circumstances which reduce some of the effects appreciably and facilitate the study of the remaining effects.

The main influence of a magnetic field in the sun is that it adds a new preferential direction to the medium which does not necessarily coincide with the direction of gravity or with the direction of wave propagation. If the direction of the magnetic field coincides with the propagation direction of the longitudinal acoustic waves there are no electromagnetic effects on the solar material and the one dimensional theory of shock wave propagation remains essentially valid. However, this is in general not the case, especially not in regions of moderate magnetic fields. In the turbulent convective layers the magnetic field is probably also in turbulent motion. It has been shown by Kulsrud (1955) that turbulent magnetic fields may increase the efficiency of the quadrupole radiation of acoustic noise generated by the compressible turbulence. On the other hand organized magnetic fields (usually stronger than the turbulent magnetic fields) which are present over large areas of active regions and in the photosphere underneath the chromospheric network, may give rise to an additional dipole emission of the acoustic noise (Osterbrock, 1961).

In a very strong field transversal motions are largely suppressed. The remaining longitudinal motions generate essentially a monopole acoustic radiation (Parker, 1964). A strong channelling of the waves along the lines of force was found by Lüst and Scholer (1966). The efficiency of the acoustic power generated is proportional to M^{2n+1}, where M is the Mach number of the turbulent

motions and n = 0, 1, 2 for monopole, dipole and quadrupole radiation. From all this we are led to the conclusion that magnetic fields will in general increase the output of acoustic waves.

The equivalent of an acoustic wave in the presence of a magnetic field is the fast mode (m.h.d.) wave. Perpendicular to the field the velocity of propagation of these waves is given by $c_f = (c_s^2 + c_a^2)^{\frac{1}{2}}$. Since the Alfvén velocity c_a increases rapidly with height above the magnetic regions the fast mode waves are strongly refracted. Hence they are probably not able to deposit the mechanical energy at the proper height in the corona (Osterbrock, 1961). However, in the form of fast mode shock waves they certainly do play a role in the heating of the chromosphere, especially above active regions. But since the propagation velocity is larger than that of gasdynamic shock waves the damping is less than for ordinary shocks. This is in agreement with the result of Uchida (1963) who showed that in a horizontal magnetic field the region of main energy dissipation is shifted towards higher layers as compared to the non-magnetic case. In a very strong magnetic field the propagation velocity is almost equal to the Alfvén velocity. Under the assumption that the energy flux and the magnetic field do not change very much over a distance of one density scale height it can be shown that the Mach number of the wavemotions is approximately proportional to $\rho^{\frac{1}{4}}$. This means that in this particular case shock waves are not to be expected, since the material velocity decreases relative to the propagation velocity when the density decreases. Consequently there is not much dissipation.

This demonstrates that in a region of strong magnetic field where more acoustic waves are generated we do not necessarily have a greater amount of heating. Although Billings (1963) showed how difficult it is to decide from the observations about the predominance of one or more of the many possible wave modes there

is good hope that future extraterrestrial observations may solve this problem.

A more important consequence of a magnetic field in a plasma with a temperature gradient is that the heat flow by thermal conduction is strongly suppressed in the direction perpendicular to the magnetic field lines. Along the field lines thermal conduction is unaffected. Perpendicular to the field lines the coefficient of heat conduction, $\varkappa_{\perp}$, is given by $\varkappa_{\perp} = \varkappa/(\omega\tau)^2$, where $\varkappa$ is the coefficient of heat conduction parallel to the direction of the magnetic field, τ is the average time between two collissions and ω is the gyro-frequency (Spitzer, 1962; Kopp, 1968). The effect is much greater for the electron conductivity than it is for the ion conductivity. In most cases it means that the heat flow by thermal conduction is essentially channelled by the magnetic field. It also means that the heating of those layers of the outer solar atmosphere, where thermal conduction is the main energy transfer, the region II, III, V, VI in figure 6 is almost entirely determined by the geometry of the magnetic field. Actually this, among others, may well be the reason why the visible structures in the outer corona appear to reflect the magnetic field geometry.

Magnetic fields in general also influence the support of the atmosphere. In a purely horizontal field the critical frequency for fast mode waves running perpendicular to the field is given by $\omega_1 = \gamma g/(c^2 + c_a^2)^{\frac{1}{2}}$ (Kuperus, 1965). Since a decrease of ω_1 results in the transmission of a greater part of the acoustic spectrum and since more energy will be generated in the magnetically active regions we are led to the conclusion that more energy is dissipated in the layers overlying magnetic regions than in the quiet regions. Consequently more energy is radiated there and the solar wind is locally enhanced. However, in certain parts of a coronal centre of activity the solar wind is almost completely suppressed by closed magnetic field structures like the observed dome shaped helmet structures. If the density in these regions is sufficiently

high radiative losses dispose of the energy deposited by the shock waves. But if on the other hand the heating exceeds the radiative losses the temperature increases locally until the domed helmet expands into a cusped helmet and eventually transforms into an open streamer along which the excess heat can be transported by the solar wind (Pneuman, 1968 a, b).

It has been shown that a necessary condition for the origin of the long lasting coronal "Condensations" is a local increase of the amount of heating (Unsöld, 1960; Pikel'ner, 1961; Kuperus, 1965).

Finally we would like to mention the possibility that part of the mechanical energy is transferred by Alfvén waves instead of acoustic waves. It was suggested by Alfvén (1947) that the corona could be heated by the dissipation of Alfvén waves due to Joule heating caused by the finite electrical conductivity. However, Cowling (1953) showed that except for hydromagnetic turbulence Joule heating is not sufficient to maintain the outer layers.

Later Piddington (1956) showed that besides ohmic losses Alfvén waves are strongly damped by the friction between the neutral and the charged constituents in a partially ionized plasma. The damping is optimal in the chromosphere at the levels where hydrogen and helium are partially ionized. In the transition region and in the corona ambipolar diffusion is not important.

In section 2 we have argued why Alfvén waves are not likely to be generated in the photosphere outside sunspot regions. The chromosphere however is more conducive for the generation of Alfvén waves. They could well be the result of either the interaction of acoustic waves or fast-mode m.h.d. waves with the magnetoplasma or of the mutual interaction of waves (Osterbrock, 1961; Pikel'ner and Lifshitz, 1964). An example of the first category is the transfer of a gravity wave into an Alfvén wave as described in section 2 (Lighthill, 1967).

7. The structure of the transition region

In the transition region between the chromosphere and the corona (the regions II, III and IV in figure 1) the physical conditions change enormously over a relatively short distance. Extremely high temperature gradients of the order of 10^4 $^oK/km$ (Kuperus, 1965) at the level where $T \approx 10^5$ ^oK are a direct consequence of the nature of heat conduction. In order to maintain energy equilibrium an increasing temperature gradient is required when the temperature decreases. Hence the chromosphere-corona interface probably has indeed a thickness of only a few tens of kilometers. Actually the analysis of the resonance lines in the far ultra-violet spectrum confirms the hypothesis that thermal conduction essentially determines the structure of the transition region at levels over $T = 10^5$ ^oK (Athay, 1966). Below this temperature the interpretation of EUV data is far too unreliable to draw this conclusion. On the other hand heat conduction is rather abruptly stopped at lower temperatures present in the chromosphere. If one then tries to give an explanation as to why the temperature in the outer solar atmosphere rises suddenly by two orders of magnitude over a very short distance one encounters a discrepancy which has been discussed by Kuperus and Athay (1967) in a qualitative way. First there is the chromosphere, which is a fairly thin region efficiently radiating away all the dissipated mechanical energy. It is now generally accepted that at a certain height the density becomes so low that at chromospheric temperatures the atmosphere cannot anymore dispose of the heat in this way. Therefore the temperature has to increase and the atmosphere adjusts to a higher temperature configuration in which again radiation from multiply ionized elements is able to dispose of the energy (Thomas and Athay, 1961; Miyamoto, 1965).

There would be no need for a corona if the atmosphere could dispose of the dissipated energy by radiating in a lower temperature

configuration. If there would be no heat conduction the temperature should rise abruptly. However, heat conduction smoothes the temperature jump and is the main cause of the finite thickness of the transition region. On the one hand the top of the chromosphere is situated at a temperature so that radiation sufficiently balances the mechanical heating, but on the other hand an additional amount of heat is dumped in those layers due to the thermal conduction from the corona backwards into the chromosphere. This conductive flux in the transition layer is about $F_c = 3 \times 10^5$ erg/cm^{2}sec. Actually this is of the same order as the radiation loss in the Lyman α line of hydrogen, which constitutes the main loss in those layers. However, this flux is originating in a shell which is assumed to be a few orders of magnitude thicker than the lowermost parts of the transition layer. It therefore seems as if the upper parts of the chromosphere and the lower parts of the transition region are in fact overheated.

Consequently Kuperus and Athay were led to the conclusion that these layers cannot be in hydrostatic equilibrium and that only the outward convection of energy associated with unstable motions can solve the enigma of the energy balance. They identified these motions with the chromospheric spicules. This suggestion about the origin of spicules differs greatly from other theories of the origin of spicules (Thomas, 1948; Osterbrock, 1961; Uchida, 1964; Parker, 1964; Wentzel and Solinger, 1967) because it essentially relates the spicules to the heating of the solar corona. Since earlier theories have always been searching for a source of spicule formation located underneath the level of spicule observations we pay here a little more attention to the "overheating hypothesis".

The influence of a magnetic field on the structure of the transition layer and the formation of spicules has been discussed by Kopp and Kuperus (1968). The main effect of a magnetic field on the structure of a region whose energy balance is almost entirely

determined by thermal conduction is the strong reduction of of conduction perpendicular to the direction of the magnetic field as was discussed in section 7. This essentially results in a channelling of the energy flow along the magnetic field. Under these conditions the flux of heat conduction is given by

$$\vec{F}_c(\ell) = -\varkappa T(\ell)^{5/2} \frac{\partial T(\ell)}{\partial \ell} \frac{\vec{B}(\ell)}{|B(\ell)|}, \qquad (15)$$

where ℓ is a coordinate measuring distance along a field line and $\vec{B}(\ell)$ is the field vector. This results in a concentration of the heat conducted backwards from the corona in the regions where the magnetic field is most concentrated. Observations show that these magnetic regions are situated at the border of the chromospheric network which is clearly visible in the Calcium H and K lines. At the same time these regions are the breeding grounds for rosettes of spicules Beckers (1967). Due to the magnetic channelling of the conducted heat flux the temperature gradient in the transition region is much larger in the spicule regions than it is in the so called interspicule regions where the magnetic field lines diverge. Therefore in the interspicule regions the heat flux decreases as one penetrates deeper into the atmosphere and radiative losses may well account for the energy balance in a hydrostatic configuration. There would be no need for motions in the interspicule regions to establish energy equilibrium in a quasi steady state. This mechanism automatically places the spicule type motions at the points where they are observed.

It would also explain approximately the measured intensities of the ultraviolet emission lines which originate in layers with temperature below 10^5 ^oK since most of these lines would of course be formed in the interspicule regions where the temperature gradient is less steep than in the spicule regions. To illustrate this figure 8 schematically shows a typical temperature structure of the transition region and inner corona, which might occur for a potential magnetic

field due to a distribution of magnetic monopoles situated in the low chromosphere. It was assumed that above a level $z = z^*$ the field is inhomogeneous. The conductive energy flux at this height was taken $F_c = 6 \times 10^5$ ergs/cm^2 sec.

For the height of the "top" of the chromosphere, which is almost equal to the height of the 10^5 ^oK isotherm, the values of 1500 km in the spicule regions and of 500 km in the interspicule regions were taken rather arbitrarily. The difference in the temperature structure of the magnetic (spicule) and "non-magnetic" (interspicule) regions is striking.

From this it is clear that it is virtually impossible to make a detailed analysis of the complicated structure of the chromosphere-corona interface without knowing the magnetic field structure. We therefore hope that high resolution ultra violet spectra and magnetic data at different heights in the atmosphere will soon become available because they are indispensable for the solution of the problem of the structure and energy balance of the outer solar atmosphere.

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References

Alfvén, H. : 1941, Ark. för Mat., Astr. o Fys. 27A no. 25.

Alfvén, H. : 1947, Mon. Not. Roy. Astron. Soc. 107, 221.

d'Angelo, N. : 1968, Astrophysic. J. 154, 401.

Athay, R.G. : 1965, The Solar Spectrum 151, ed. C. de Jager, D. Reidel, Dordrecht.

Athay, R.G. : 1966a, Astrophys. J. 145, 784.

Athay, R.G. : 1966b, Astrophys. J. 146, 223.

Barnes, A. : 1968, Astrophys. J. 154, 751.

Beckers, J.M. : 1968, Solar Physics 3, 367.

Biermann, L. : 1946, Naturwiss. 33, 118.

Biermann, L. : 1948, Z.f.Astroph. 25, 161.

Biermann, L. : 1966, "The Fine Structure of the Solar Atmosphere", Colloq. held at the Observatory Anacapri of the Fraunhofer Inst., ed. K.O. Kiepenheuer, page 116. (Franz Steiner Verlag - Wiesbaden).

Biermann, L. and ten Bruggencate, P.: 1947, Göttingen, Univ. Sternw. Veröff. 83.

Billings, D.E.: 1963, Astrophys. J. 137, 592.

Billings, D.E.: 1966, A Guide to the Solar Corona, Academic Press, New York.

Bird, G.A. : 1964a, Astrophys. J. 139, 675.

Bird, G.A. : 1964b, Astrophys. J. 139, 684.

Bird, G.A. : 1964c, Astrophys. J. 140, 288.

Bird, G.A. : 1965a, Astrophys. J. 141, 1455.

Bird, G.A. : 1965b, Phys. of Fluids 8, 2289.

Bodenheimer, P., Brandt, J.C. and Robbins, R.R.: 1963, Icarus 2, 411.

Bondi, H., Hoyle, F. and Littleton, R.A.: 1947, Mon. Not. Roy. Astron. Soc. 107, 184.

Brinkley, S.R. and Kirkwood, J.G.: 1947, Phys. Review 71, 606.

Burgess, A. : 1964, Astrophys. J. 139, 776.

Chapman, S. : 1954, Astrophys. J. 120, 151.

Chapman, S. : 1957, Smithson. Contr. Astrophys. 2, 1.

Chisnell, R.E.: 1955, Proc. Roy. Soc. London A 232, 350.

Cowling, T.G. : 1953, "The Sun", ed. G.P. Kuiper, 579-583; (The University of Chicago Press, Chicago).

Dubov, E.E. : 1960a, Izvestia Krimsk. Astrof. Obs. 22, 101. translation: AFCRL - 63 - 265.

Dubov, E.E. : 1960b, Izvestia Krimsk. Astrof. Obs. 23, 362.

Dubov, E.E. : 1965, Astron. Zhurnal 42, 1014. translation: 1966, Soviet Astron. 9, 782.

Eckart, C. : 1960, Hydrodyn. of Oceans and Atmosph., (New York: Pergamon).

Edlén, B. : 1943, Z. f. Astrophys. 22, 30.

Evans, J.W. and Michard, R.: 1962, Astrophys. J. 136, 493.

Ffowcs Williams, J.E.: 1963, Phil. Trans. Roy. Soc. London A 255, 469.

Giovanelli, R.G. : 1949, Mon. Not. Roy. Astron. Soc. 109, 372.

Hines, C.O. : 1960, Can. Journ. of Phys. 38, 1441.

van de Hulst, H.C. : 1949, Probl. of Cosm. Aerodyn. 45, ed. J.M. Burgers and J.H. Oort (Central Air Doc. Off., Dayton 2, Ohio)

van de Hulst, H.C. : 1953, "The Sun", page 207, ed. G.P. Kuiper Chicago Univ. Press.

de Jager, C. and Kuperus, M.: 1961, Bull. Astron. Neth. 16, 71.

Kato, S. : 1966, Astrophys. J. 143, 372.

Kopp, R.A. : 1968, Thesis Harvard University Scientific Report No. 4, AFCRL - 68 - 0312.

Kopp, R.A. and Kuperus, M.: 1968, Solar Phys. 4, 212.

Kulsrud, R.M. : 1955, Astrophys. J. 121, 461.

Kuperus, M. : 1964, Joint Discussion on Aerodyn. Phen. in stellar atmosph., Transactions I.A.U. 12B, 564.

Kuperus, M. : 1965, Rech. Astron. Obs. Utrecht 17, 1.

Kuperus, M. : 1969, in preparation.

Kuperus, M. and Athay, R.G.: 1967, Solar Phys. 1, 361.

Landau, L. and Lifshitz, E.M.: 1959, Fluid Mechanics. (New York: Pergamon)

Leighton, R.B. : 1964, Ann. Rev. Astron. and Astrophys. 1, 19.

Lighthill, M.J. : 1952, Proc. Roy. Soc. London A 211, 564.

Lighthill, M.J. : 1954, Proc. Roy. Soc. London A 222, 1.

Lighthill, M.J. : 1960, Phil. Transact. Roy. Soc. A 252, 397.

Lighthill, M.J. : 1962, Proc. Roy. Soc. London A 267, 147.

Lighthill, M.J. : 1967, "Aerodynamic Phenomena in Stellar Atmospheres", ed. R.N. Thomas, I.A.U. Symp. No. 28, 429.

Lüst, R. and Scholer, M.: 1966, Z.f.Naturforsch. 21, 1098.

Lyot, B. : 1932, Zs. f. Astrophys. 5, 73.

Lyot, B. : 1937, L'Astronomie 51, 203.

Mäckle, R. : 1968, Mitt. Astron. Gesellsch. 25, 200.

Miyamoto, S. : 1965, Publ. Astron. Soc. Japan 17, 127.

Moore, D.W. and Spiegel, E.A.: 1964, Astrophys. J. 139, 48.

Moore, D.W. : 1967, "Aerodynamic Phenomena in stellar atmospheres ed. R.N. Thomas, I.A.U. Symp. 28, 429.

Newkirk, G. : 1967, Ann. Rev. Astron. and Astrophys. 5, 213.

Noble, L.M. and Scarf, F.L.: 1963, Astrophys. J. 138, 1169.

Ôno, Y., Sakashita, S. and Ohyama, H.: 1961, Progr. Theor. Phys. Kyoto Suppl. 20, 85.

Orrall, F.Q. and Zirker, J.B.: 1961, Astrophys. J. 134, 72.

Oster, L. : 1956, Z. f. Astrophys. 40, 28.

Osterbrock, D.E. : 1961, Astrophys. J. 134, 347.

Pagel, B.E.J. : 1964, Ann. Rev. Astron. and Astrophys. 2, 267.

Parker, E.N. : 1958, Astrophys. J. 128, 677.

Parker, E.N. : 1963, "Interplanetary Dynamical Processes" (New York: Interscience publishers)

Parker, E.N. : 1964, Astrophys. J. 140, 1170.

Parker, E.N. : 1965, Space Sci. Rev. 4, 666.

Pikel'ner, S.B. : 1950, Izv. Krymsk. Astrofiz. Obs. 5, 34.

Pikel'ner, S.B. : 1961, Astron. Zhurn. 38, 552.
translation: Soviet Astron. 1961, 5, 412.

Pikel'ner, S.B. and Livshitz, M.A.: 1964, Astron. Zhu. 41, 1007.
1965, Soviet Astron. 8, 808.

Piddington, J.H.: 1956, Mon. Not. Roy. Astron. Soc. 116, 314.

Pneuman, G.W. : 1968a, Solar Phys. 4, 578.

Pneuman, G.W. : 1968b, Solar Phys. 5, 595.

Proudman, I. : 1952, Proc. Roy. Soc. London A 214, 119.

Saito, M. : 1964, Publ. Astron. Soc. Japan 16, 179.

Scarf, F.L. and Noble, L.M.: 1965, Astrophys. J. 141, 1479.

Schatzman, E. : 1949, Ann. d'Astrophys. 12, 203.

Schatzman, E. : 1967, "Solar Physics" proceedings of NATO Summer course held in Lagonissi, Greece, ed. J.N. Xanthakis, 133.

Schatzman, E. and Souffrin, P.: 1967, Ann. Rev. Astron. and Astrophys. 5, 67.

Schirmer, H. : 1950, Z. f. Astrophys. 27, 132.

Schwarzschild, M.: 1948, Astrophys. J. 107, 1.

Shklovskii, I.S.: 1944, "thesis, Moscow State Univ. (Moscow)".

Shklovskii, I.S.: 1965, "Physics of the Solar Corona" (Pergamon Press)

Souffrin, P. : 1966, Ann. d'Astrophys. 29, 55.

Stein, R.F. : 1967, Solar Physics 2, 385.

Sturrock, P.A. : 1964, Nature 203, 285.

Spitzer, L. : 1962, "Physics of fully ionized gases" (New York: Interscience publishers).

Thomas, R.N. : 1948, Astrophys. J. 108, 130.

Thomas, R.N. and Athay, R.G.: 1961, "Physics of the solar chromosphere" (New York, Interscience publ.).

Tolstoy, I. : 1963, Rev. Mod. Phys. 35, 207.

Uchida, Y. : 1961, Publ. Astron. Soc. Japan 13, 321.

Uchida, Y. : 1963, Publ. Astron. Soc. Japan 15, 376.

Uchida, Y. : 1965, Astrophys. J. 142, 335.

Uchida, Y. : 1967, "Aerodynamic Phenomena in stellar atmospheres" ed. R.N. Thomas, I.A.U. Symp. 28, 456.

Ulmschneider, P.: 1967, Z. f. Astrophys. 67, 193.

Unno, W. : 1964, Joint discussion on aerodyn. phen. in stellar atmospheres, Transactions I.A.U. 12B, 555.

Unno, W. and Kato, S.: 1962, Publ. Astron. Soc. Japan 14, 417.

Unsöld, A. : 1960, Z. f. Astrophys. 50, 57.

Vitense, E. : 1953, Z. f. Astrophys. 32, 135.

Wentzel, D.G. and Solinger, A.B.: 1967, Astrophys. J. 148, 877.

Weymann, R. : 1960, Astrophys. J. 132, 452.

Whitaker, W.A. : 1963, Astrophys. J. 137, 914.

Whitham, G.B. : 1958, J. Fluid Mech. 4, 337.

Whitney, C.A. : 1957, Smithson. Contr. to Astrophys. 2, No. 12.

Woolley, R.v.d.R. and Allen, M.N.: 1950, Mon. Not. Roy. Astr. Soc. 110, 358.

Zel'dovich, Ya.B. and Raizer, Yu.P.: 1966, "Physics of Shock waves and High-Temperature hydrodynamic Phenomena", vol. I, 72.
(Academic Press)

Captions to the figures.

Figure 1: Propagation surfaces for acoustic waves in a gravitational atmosphere. The arrows indicate the direction of the group velocity.

Figure 2: The acoustic power output as a function of depth in the turbulent layers of the convection zone.

Figure 3: The directional dependence of the acoustic power output in convective turbulence.

Figure 4: Deformation of a sound wave with a period of 200 sec during propagation through an isothermal atmosphere of $T = 5000\ ^{o}K$. The velocity is plotted against time for three different heights, expressed in the scale height H.

Figure 5: The dissipation of shock waves.

Figure 6: Schematic division of the outer solar atmosphere according to energy-balance considerations (after Kopp).

Figure 7: Comparison of models for the outer solar atmosphere.

Figure 8: A schematic cross-section of the region upper chromosphere-lower corona showing magnetic field lines, isotherms, and directions of the observed motions.

Figure 1

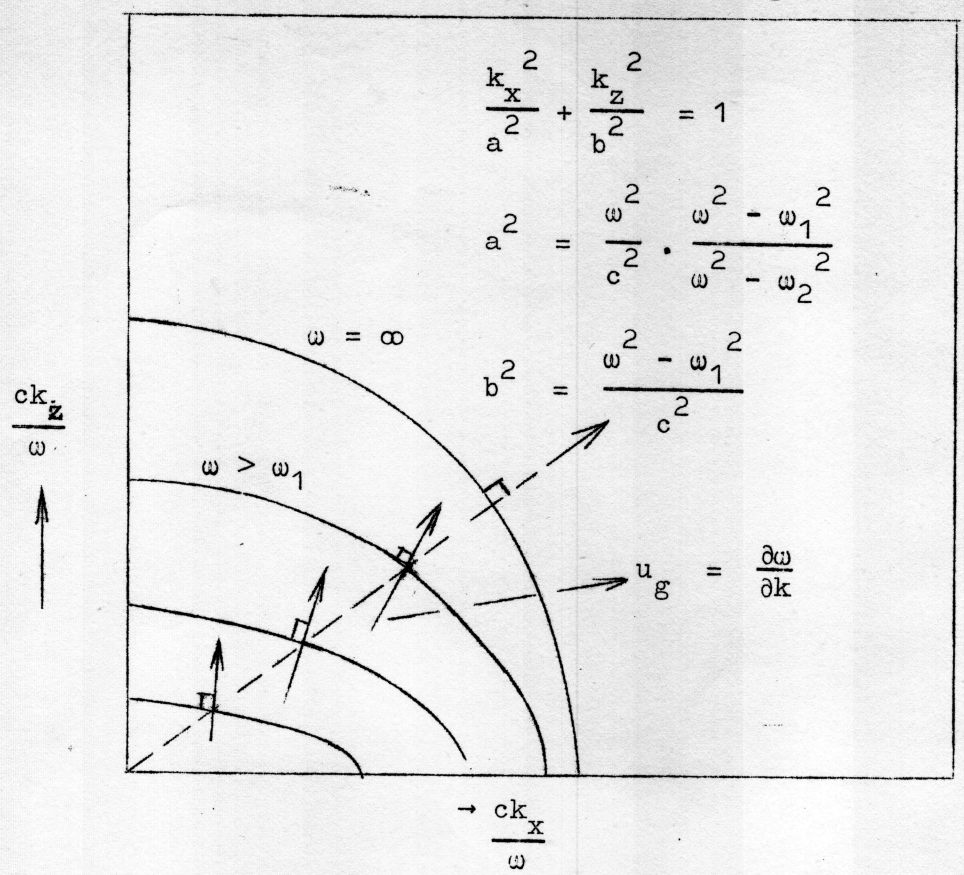


Figure 2

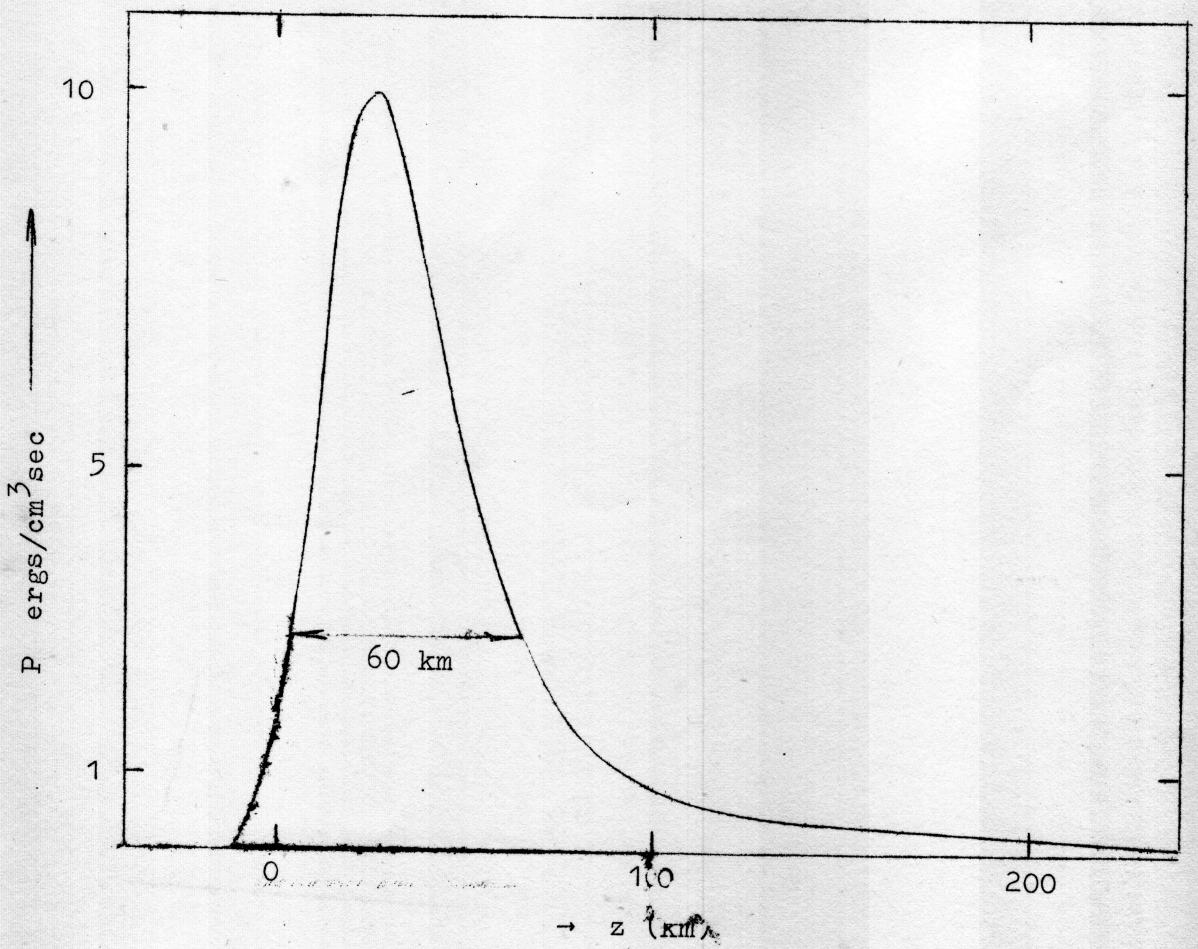


Figure 3

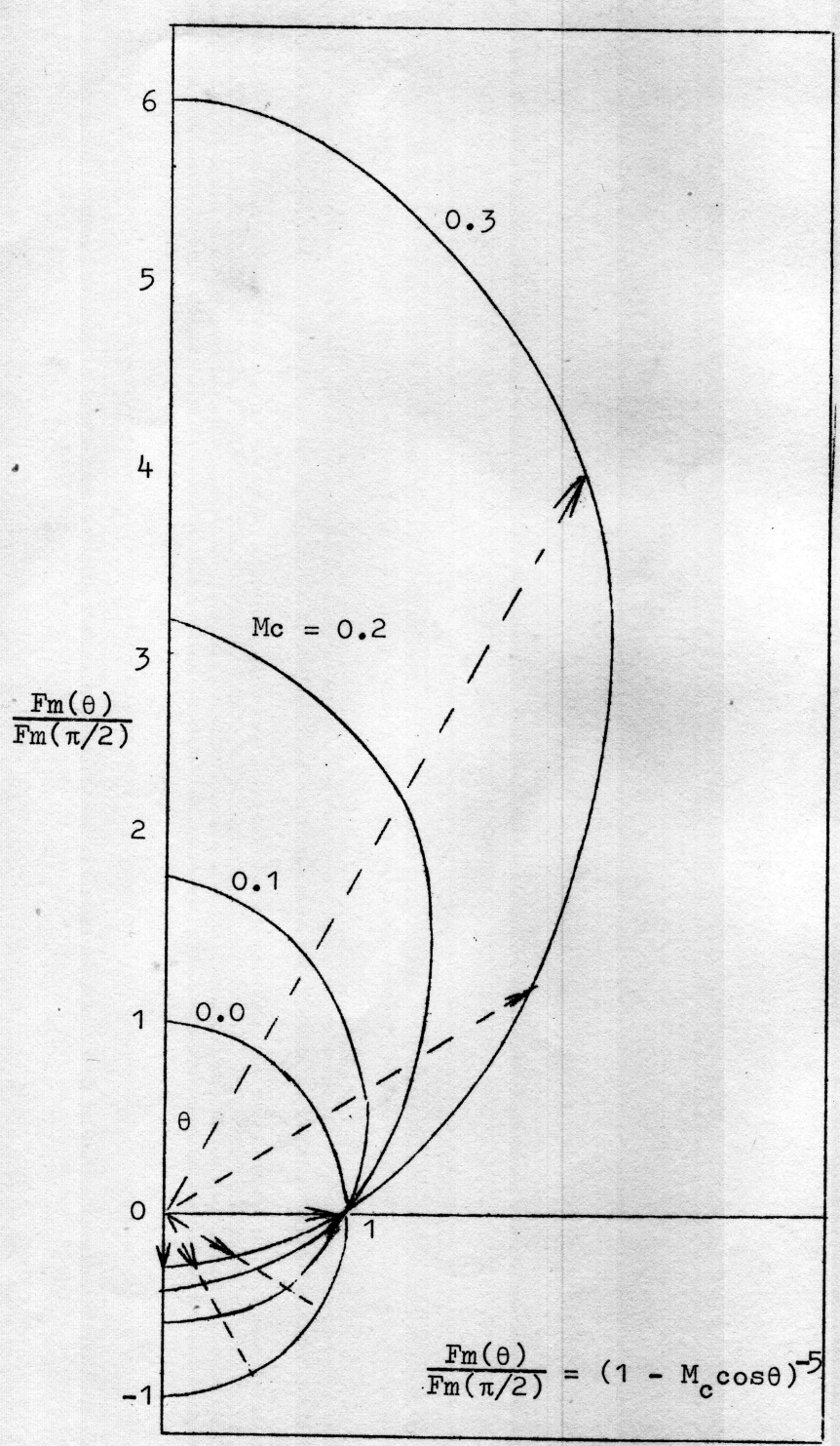
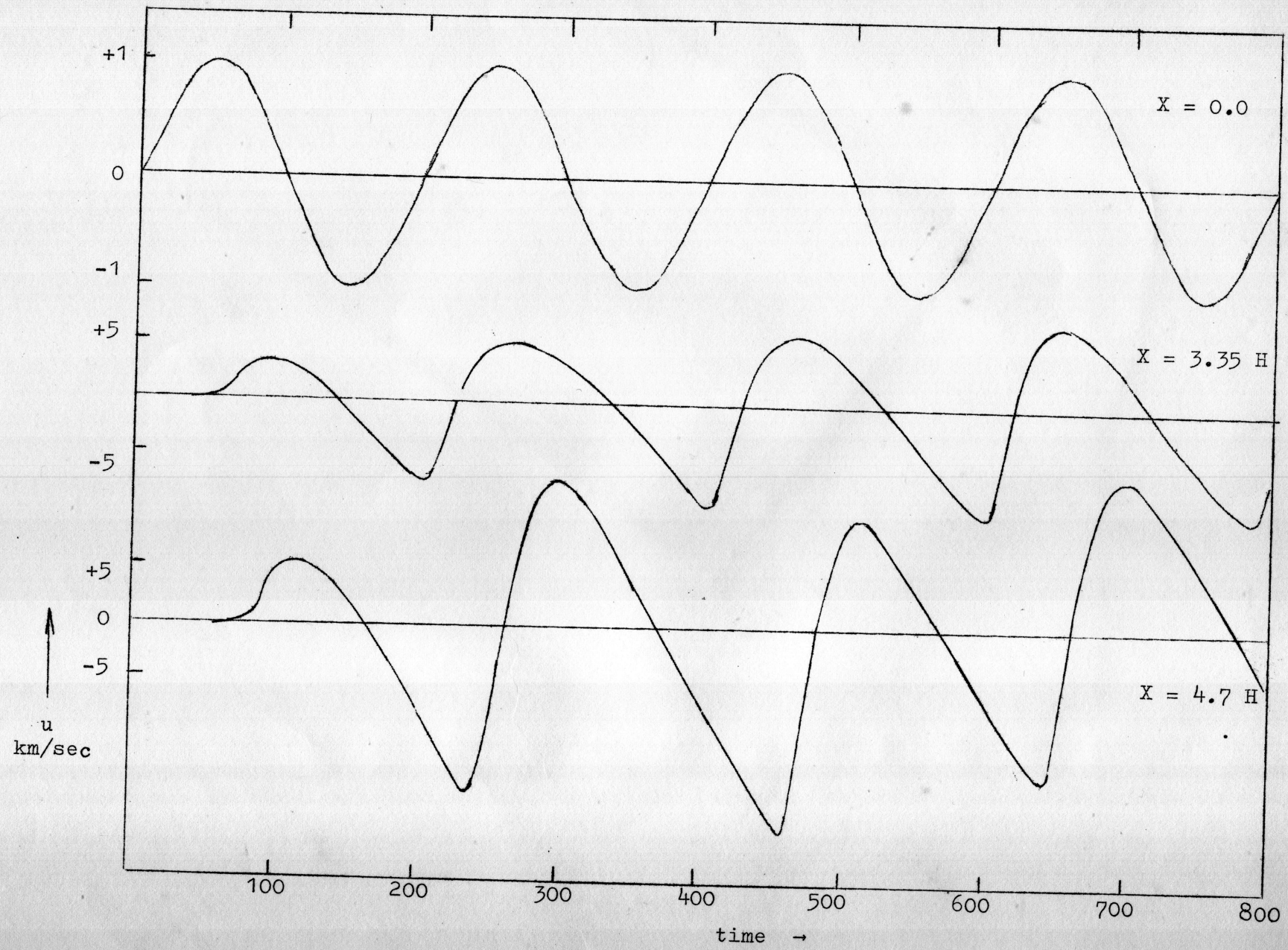


Figure 4


+1
0
-1
+5
-5
+5
0
-5
X = 0.0
X = 3.35 H
X = 4.7 H
u
km/sec
100
200
300
400
500
600
700
800
time →

Sound wave: ξ vertical

$$\rho_0 \frac{\partial^2 \xi}{\partial t^2} = -\nabla p_1 - \rho_1 g \quad \otimes$$

$$\frac{\partial \rho_1}{\partial t} + (\vec{v}\cdot\vec{\nabla})\rho_0 + \rho_0 \nabla\cdot\vec{v} = 0 \Rightarrow \rho_1 = -\xi \frac{d\rho_0}{dz} - \rho_0 \frac{d\xi}{dz} \quad *$$

adiab.: $\frac{dP}{dt} = \frac{\partial P}{\rho}\frac{d\rho}{dt}$

$$\frac{\partial \rho_1}{\partial t} + (\vec{v}\cdot\nabla)\rho_0 = c^2 \frac{\partial \rho_1}{\partial t} + c^2(\vec{v}\cdot\nabla)\rho_0$$

$$p_1 = c^2\rho_1 + c^2\xi\frac{d\rho_0}{dz} - \xi\frac{dp_0}{dz}$$

$$\overset{*}{=} c^2\left(-\xi\frac{d\rho_0}{dz} - \rho_0\frac{d\xi}{dz}\right) + c^2\xi\frac{d\rho_0}{dz} - \xi\frac{dp_0}{dz}$$

$$= -\rho_0 c^2\frac{d\xi}{dz} + \rho_0 g\xi$$

$$\frac{dp_0}{dz} = -\rho_0 g$$

$$\otimes \quad \frac{\partial^2\xi}{\partial t^2} = +c^2\frac{d^2\xi}{dz^2} + \frac{c^2}{\rho_0}\frac{d\rho_0}{dz}\frac{d\xi}{dz} + \xi g\frac{d\rho_0}{dz} + \rho_0 g\frac{d\xi}{dz}$$

2.1 Trial: $H = H + h$; $v = v$; $\rho = \rho_0 + \rho$: tot op 2e orde:

$$i\frac{\partial v}{\partial t} - \frac{i}{4\pi\rho_0}[\nabla\times h, H] + i(v\cdot\nabla)v - \frac{i}{4\pi\rho_0}[\nabla\times h, h] = 0$$

$$i\frac{\partial h}{\partial t} - i\nabla\times[v, H] + \alpha i\nabla\times\frac{1}{\rho_0}[\nabla\times h, H] - i\nabla\times[v, h] + \alpha i\nabla\frac{1}{\rho_0}[\nabla\times h, h] = 0$$

$$i\frac{\partial \rho}{\partial t} + i\nabla(\rho_0 v) + i\nabla\cdot(\rho v) = 0$$

$$\varphi = \begin{pmatrix} v \\ h \\ \rho \end{pmatrix}$$

$$i\frac{\partial \varphi}{\partial t} + \hat{H}_0\varphi + \hat{H}_1\{\varphi, \varphi\} = 0$$

$$\hat{H}_0 = \begin{pmatrix} 0 & +\frac{i}{4\pi\rho_0} H\times(\nabla\times\ ) & 0 \\ i\nabla\times(H\times\ ) & -\frac{\alpha i}{\rho_0}\nabla\times(H\times(\nabla\times\ )) & 0 \\ i\rho_0\nabla\cdot & 0 & 0 \end{pmatrix}$$

S STRUCTURE & physical properties of chromosph. & corona

O observational indications for heating by mechanical waves

E energy losses